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An Introduction to Quark-Gluon Plasma and High Energy Heavy Ion Collisions [1]

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Abstract

The quark-gluon plasma, and how it might be produced in ultra-relativistic nuclear collisions is reviewed. I briefly introduce the quark-gluon plasma, and what we might learn from studying it. I then discuss what has been learned from the recent results from the CERN oxygen run. I then attempt to address the issue of whether A = 16 and E = 200 Gev are sufficient to make a quark-gluon plasma. I discuss strangeness and charm production as well as electromagnetic probes of the plasma, but do not discuss hydrodynamics, multiplicity and Et distributions.

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1 Introduction

Perhaps the most interesting aspect of high energy nuclear interactions is that it may allow for tests of unique features of QCD. These features reflect non-perturbative phenomenon such as confinement and chiral symmetry breaking. In this talk I shall first give an overview of current theoretical understanding of these non-perturbative phenomenon.

To study matter at densities of the order of and larger than those typical of QCD, we must study either the collisions of ultra-relativistic nuclei, or very high multiplicity fluctuations in hadron-hadron collisions. We shall see that simple arguments suggest that densities far in excess of those typical of ordinary nuclei may be achieved under such extreme conditions. I will later briefly discuss a few suggested experimental probes of high density matter as it might be produced in such collisions.

In this section I shall discuss the properties of hadronic matter at high energy density. The word high implies a scale for the measurement of the energy density. Such a scale may be provided by a variety of estimates, all of which agree on the order of magnitude of a typical density scale for hadronic matter. The first is the energy density of nuclear matter. With m the proton mass, R_A the nuclear radius, and A the nuclear baryon number, the density of nuclear matter is

$$\rho_A \sim \frac{Am}{\frac{4}{3}\pi R_A^3} \sim .14\ Gev/Fm^3 \qquad (1)$$

We can also use Eq. 1 to estimate the energy density inside a proton. If we use a proton radius of .8 Fm, Eq. 1 gives

$$\rho_p \sim .5\ Gev/Fm^3 \qquad (2)$$

There is a good deal of uncertainty in this estimate of ρ_p. We might have instead used the MIT bag radius, or a proton hard core radius, corresponding to an order of magnitude uncertainty in Eq. 2. Finally, another estimate comes from dimensional grounds using the value of the QCD Λ parameter, suitably defined as Λ_{ms} or Λ_{mom}, as the dimensional scale factor. Using the Λ parameter, we find

$$\rho_{QCD} \sim \Lambda^4 \sim .2\ Gev/Fm^3 \qquad (3)$$

Again there is an order of magnitude uncertainty both due to the lack of precise experimental knowledge of Λ, and differences induced by using alternative sensible definitions of Λ.

In all of the above energy density estimates, the typical scale was in the range of several hundreds of Mev/Fm^3 to several Gev/Fm^3. At energy densities low compared to this scale, we presumably have a low density gas of the ordinary constituents of hadronic matter, that is, mesons and nucleons. At densities very high compared to this scale, we expect an asymptotically free gas of quarks and gluons.[77] At intermediate energy densities, we expect that the properties of matter will interpolate between these dramatically different phases of matter. There may or may not be true phase changes at some intermediate densities.

To understand how such a transition might come about, consider the example of QCD in the limit of a large number of colors, N_C.[77] Recall that extensive quantities such as the energy density, ϵ, or entropy density, σ, measure the number of degrees of freedom of a system. The dimensionless quantities ϵ/T^4 or σ/T^3 should be of the order of the number of degrees of freedom. For hadronic matter, the number of degrees of freedom relevant at low density are the number of low mass hadrons. Since matter is confined at low density, the number of such degrees of freedom is $N_{dof} \sim 1$ in terms of the number of colors. At high energy density, the relevant number of degrees of freedom are those of unconfined quarks and gluons. The gluons dominate and give $N_{dof} \sim N_C^2$. Therefore in the large N limit, the number of degrees of freedom change by an infinite amount.

Assuming that the transition occurs at finite temperature in the large N_C limit, as is verified by Monte-Carlo simulation, this result can be interpreted in two ways.[77] From the vantage point of a high density world of gluons, the asymptotic energy density is finite, but at low energy density at some finite temperature the energy density goes to zero. The energy density itself is therefore an order parameter for a phase transition, and there is a limiting lowest temperature. Viewed from the low density hadronic world, there is some limiting temperature where the energy density and entropy density become infinite. Here there is a Hagedorn limiting temperature.[77]

For $N_C = 3$, the above statements are only approximate. The number of degrees of freedom of low mass mesons is

$$N_{dof} \sim N_F^2 \sim 4 \tag{4}$$

where we have taken the number of low mass quarks to be $N_F \sim 2$ for the up and down quarks. The number of degrees of freedom of a quark-gluon plasma is on the other hand

$$N_{dof} \sim 40 \tag{5}$$

The number of degrees of freedom might change in a narrow temperature range, or there might be a true phase transition where the degrees of freedom change by an order of magnitude, if our speculations concerning the large N_C limit are applicable.

Results of a Monte-Carlo simulation of the energy density are shown in Fig. 1.[??,??] These results are typical of the qualitative results arising from lattice Monte-Carlo simulation. The precise values of the energy density are difficult to estimate as is the scale for the temperature. The figure does make clear the essential point, on which all Monte-Carlo simulations agree, that the number of degrees of freedom of hadronic matter changes by an order of magnitude in a narrowly defined range of temperature. There is apparently a first order phase transition for SU(3) Yang-Mills theory in the absence of fermions, and a rapid transition which may or may not be a first order transition for SU(3) Yang-Mills theory with two or three flavors of massless quarks.

For Yang-Mills theory in the absence of dynamical quarks, there is a local order parameter which probes the confinement or deconfinement of a system. This order parameter measures the exponential of the free energy difference between the thermal system with and without the presence of a single static test quark inserted as a probe,

$$< L > = e^{-\beta F_q} \tag{6}$$

As originally proposed by Polyakov[??] and Susskind,[??] and developed in Monte-Carlo studies,[??] the Polyakov loop is a Wilson loop at the position of the quark which evolves only in time and is closed by virtue of the thermal boundary conditions which make the system have a finite extent in Euclidian time. The two phases of the theory are the confined and unconfined phases where

$$e^{-\beta F_q} \sim \; finite \; if \; confined, \; or \; 0 \; if \; deconfined \tag{7}$$

This quantity is an order parameter for a confinement-deconfinement in theories without fermions or in the large N_C limit in theories with fermions (in the fundamental representation of the gauge group). If there are fermions in the fundamental representation, in the 'confined phase' dynamical fermions may form a bound state with a heavy test quark, so the free energy is finite in what would be the confined phase.[77] Since it is already finite in the deconfined phase, the free energy of a static test quark does not provide an order parameter.

Although $< L >$ is not an order parameter, Monte-Carlo simulations with dynamical fermions show that $< L >$ changes very rapidly in a narrow range of temperatures.[77] For SU(3) lattice gauge theory without dynamical quarks, when $< L >$ is a true order parameter, there is a noticeable discontinuous change. It is not entirely clear whether there is a discontinuous change corresponding to a true phase change for the theory with fermions.

In the limit of large dynamical quark mass the quarks are no longer important at any finite temperature and decouple. In this limit the confinement-deconfinement phase transitions is a well defined concept with an order parameter which measures a phase change. At zero quark masses there is another phase transition which may be carefully defined, that is, the chiral symmetry restoration phase transition. Chiral symmetry is a continuous global symmetry of the QCD lagrangian in the limit of zero quark mass. Its realization would require that all non-zero mass baryons have partners of degenerate mass and opposite parity. Since this is far from true for the spectrum of baryons observed in nature, chiral symmetry must be broken. Breaking the continuous global symmetry generates a massless Goldstone boson, which we identify with the light mass pion. As a consequence of the breaking of chiral symmetry, the quarks acquire dynamical masses, which may be seen by computing $< \overline{\Psi}\Psi >$. For the chiral symmetric phase, $< \overline{\Psi}\Psi > = 0$, and is non-zero in the broken phase.

For not unreasonable values of the quark masses, there appears to be a rapid change in $< \overline{\Psi}\Psi >$ at about the same place where the order parameter $< L >$ changes rapidly. We conclude therefore that chiral symmetry is approximately restored at the same temperature where quarks stop being approximately confined. The word approximately is important here since absolute confinement or absolute chiral symmetry is impossible for finite mass dynamical quarks.

We can now conjecture on the phase diagram in the temperature mass plane. It is important to realize that we may physically vary the temperature, but not the masses of quarks. Theoretically in a Monte-Carlo simulation, these masses may be changed, but they cannot be changed in nature. It is also important to realize that the mass-temperature diagram represents an over simplification to the case of equal mass quarks. With different mass quarks, the diagram has more variables and is more complicated.

To plot this diagram, we first discuss the limiting case $m = \infty$. Here there should be a first order confinement-deconfinement phase transition along the T axis. Since a discontinuous change will not be removed by a large but finite quark mass, this first order phase change must be a line of transitions in the m-T.Along the $m = 0$ axis there is a chiral symmetry restoration transition. By the arguments of Pisarski and Wilczek,[??] this transition is first order, and therefore must generate a line of transitions which extends into the m-T plane.

Of course, we do not know what happens with these two lines of transitions, whether they join or never meet, or pass through one another etc. There may be no true phase transition at the values of masses which are physically relevant, or there may be one or two which are the continuation of the chiral transition from zero mass and the confinement-deconfinement transition from infinite mass. The weight of the evidence from Monte-Carlo numerical simulation suggests a very large transition in the properties of matter in a very narrow temperature range, and not much more than that can be said at present. There are a variety of conflicting claims as to whether or not there is a true first order transition at physically relevant masses.[??]

There have been serious attempts to obtain reliable quantitative measures of the properties of matter from Monte-Carlo simulation.[??,??] The only truly reliable numbers have been extracted for the unphysical case of $N_F = 0$, that is, no dynamical fermions. It has been shown that the critical temperature of the confinement-deconfinement transition is

$$T_C = 220 \pm 50\, Mev \tag{8}$$

by fitting the potential computed in these theories and comparing it with the potential which fits charmonium. This corresponds to an energy density of $1-2\ Gev/Fm^3$ required to make a quark-gluon plasma. These results now appear to be valid for the continuum limit, and seem to be fairly good.

The numerical situation for QCD with $N_F = 2-3$ is not nearly so good. The qualitative results have been summarized above, but it is premature to draw any firm conclusions about numbers.

2 How to Make a Plasma

The collisions of ultra-relativistic nuclei and fluctuations in $\bar{p}p$ collisions provide the possibility of producing a quark-gluon plasma in a controlled experimental environment.[??,??] Such a collision is shown in Fig. 2 where two nuclei of transverse radius R collide in the center of mass frame. The longitudinal size of the nuclei is Lorentz contracted.

There is a scale implicit in the Lorentz contraction. Once the nuclei have a large enough Lorentz gamma factor so that they would be contracted to a size less than some typical hadronic length scale, possibly a fermi, the Lorentz contraction of virtual quanta with energy corresponding to this length scale stops. Below the beam energy appropriate for this gamma factor, the nuclei Lorentz contract. This energy is

$$E^0_{CM} = m\gamma = \frac{mR}{l_0} = 7-70\ Gev \qquad (9)$$

for uranium nuclei and the hadronic distance scale $l_0 \sim .1-1\ Fm$. Here and in the rest of this paper, we shall quote the center of mass energy in Gev per nucleon in each nucleus.

We expect qualitative differences in the scattering above E^0_{CM}. Another equivalent estimate of E^0_{CM} is given by estimating the energy at which the fragmentation regions of the two nuclei separate. At energies greater than E^0_{CM} there will be a central region between the two colliding nuclei, which will have small net baryon number density. This separation is shown in the dual-parton model computations of O-Pb collisions[??], shown in Fig. 3.

An important fact to remember about the matter formed in the collision of two ultra-relativistic nuclei is that it is born expanding in the longitudinal direction. This is because particles are formed with a more or less uniform density in rapidity.

Since these particles follow a trajectory which has its origin approximately at $x = t = 0$, and there is a large dispersion in particle velocities, there will be a large longitudinal velocity gradient built into the initial matter distribution. There should be no transverse expansion in the initial condition since we expect a random orientation in the transverse momentum of produced particles. It can be shown that if the distribution of produced particles is uniform in rapidity, the expansion is initially a 1+1 dimensional similarity expansion, and the density of particles decreases like 1/t.

The initial energy density may be estimated on dimensional grounds. The initial energy density should be proportional to the initial rapidity density per unit transverse area. The energy per particle should be of the order of the typical transverse momentum per particle. The longitudinal distance scale and p_T are correlated at early time by the uncertainty principle, since initially the matter appears in a quantum mechanical state, $p_T \sim 1/l_o$. We therefore have

$$\epsilon_i \sim \frac{dN}{dy} \frac{1}{\pi R^2} p_T^2|_{t=t_i} \tag{10}$$

The initial time t_i will be chosen as the earliest time we believe that the matter may be described as approximately expanding as a perfect fluid.

If the matter expands approximately as a perfect fluid, then ϵ_i may be bounded by parameters which are experimentally measured at late times after the matter decouples, that is, after the pions present in the late state of evolution of the matter have stopped scattering from one another, and are experimentally observed. We first use that the rapidity density in perfect fluid hydrodynamic expansion is proportional to the entropy and because entropy is conserved, one can prove that dN/dy is also conserved, at least in the central region.[17] Since the system cools as it expands, p_T is a monotonically decreasing function of time. (Some of the transverse momentum is recovered by transverse flow, but p_T nevertheless monotonically decreases.) We find therefore that

$$\epsilon_i > p_t^2 \frac{1}{\pi R^2} \frac{dN}{dy} \tag{11}$$

In this equation, all quantities are experimentally observable.

It should be strongly emphasized that the above estimate only applies to the central region for collisions of equal A nuclei of large A at very high energy. Therefore, the above formula does not apply for the asymmetric A, low energy collisions at CERN. Estimates of the energy density in the fragmentation region for asymmetric nuclei have not yet been attempted.

The initial energy density might be much larger than this estimate for a variety of reasons. In fluctuations in $\bar{p}p$ collisions, the multiplicity may be much larger. In nuclear collisions, the initial p_T may be much larger than is typical of the final state. This initial p_T may be determined by kinetic theory arguments, and might be in the range of $.4 - 2\ Gev$,[??,??] corresponding to uncertainty in the energy density of at least an order of magnitude. The initial transverse momentum, and correspondingly, the initial time, may even depend upon the nuclear baryon number A.[??] I think the best estimates of the achievable energy densities in central collisions of large nuclei is $2 - 200\ Gev/Fm^3$. This corresponds to an initial temperature in the range of $T_i \sim 200 - 700\ Mev$.

To achieve very high energy densities, however requires very high energy densities. If the initial formation time is $t_o = C/T$ where C is of order one , to acquire a temperature of 500 Mev would require that the nulcei be lorentz contracted to a size of .4 Fm, which for uranium requires a center of mass energy of 40 Gev/A. Also, it is quite likely that the maximum possible temperatures are only achieved for very large A[??].

To make a convincing case that there is sufficient time for the formation and evolution of a quark-gluon plasma as an approximate perfect fluid, the expansion rate of the system should be compared to a typical particle collision time. When the collision time is much less than the expansion time, the system should expand approximately adiabatically as a perfect fluid. Since entropy is conserved, the initial and final times for expansion in d dimensions are related by

$$\left(\frac{t_f}{t_i}\right)^d = \frac{N^i_{dof}}{N^f_{dof}} \frac{T_i^3}{T_f^3} \sim 10 - 10^4 \qquad (12)$$

where N_{dof} are the number of particle degrees of freedom. At early time, the expansion is 1 dimensional, and later times becomes three dimensional. We estimate therefore that $t_f/t_i \sim 10 - 10^3$. Detailed hydrodynamic computations show that the final decoupling time is probably somewhere in the range of $t_f \sim 20 - 50\ Fm/c$.[??]

Large nuclei are clearly the more favored system for producing and studying a quark-gluon plasma. This follows simply from the facts that the average energy density achieved is larger, and that the system is physically larger in transverse extent. We require $\lambda_{scat} << R_{nuc}$ in order for a perfect fluid hydrodynamic treatment to be sensible. Estimates of λ_{scat} give $.1 - 1\ Fm$.[??,??]

Experimental data exists which throws some light on the size of systems necessary for fluid dynamic effects to become important. At Bevalac energies, the flow of hadronic matter was studied in nuclear collisions.[??,??] In collisions of nuclei of small impact parameter, single particle collisions occur at large transverse momentum. The nuclei do not collectively flow in a given transverse direction unless there are subsequent rescatterings among the constituents of the nuclei. If these subsequent rescatterings do not occur, the transverse momentum of each particle is randomly oriented. To get collective flow, one needs rescattering, and this should be enhanced in collisions at small impact parameter, and collisions of large A nuclei.

In Fig. 4, the flow angle is plotted for various measures of the impact parameter (large impact parameters at the top and small at the bottom of the figure) for various nuclei (small on the left and large on the right). Little evidence of flow is shown for nuclei as large as calcium, and collective effects begin to become important for nuclei of the size of niobium.

The current experiments at CERN may allow for some determination of the energy densities which might be achievable in high energy heavy ion collisions at asymptotically high energy and for very large A. To sort this out from the data, one must have models to compare the data with. The data which is now available is primarilly for E_t and dN/dy distributions. In principle the correlation between these variables can determine whether there is thermalization. For thermal models the p_t is enhanced due to rescattering. This is a small effect for pions, but is a larger effect for nucleons, as will be discussed in later talks. In Table 1, I give a list of various models which attempt to describe nuclear collisions and the distinguishing features which may allow their resolution.

Models of Nuclear Collisions

Model	Thermalization?	p_t Enhancement?
DPM, Hi-Jet	no	no
Lund, Rope model	no	some
Nuclear Cascade	some	yes
QGP	yes	yes

3 Probes of the Quark-Gluon Plasma

In Table 2, various experimental probes of the quark-gluon plasma are presented.

Probes of the Quark-Gluon Plasma

Probe	Physics
Photons and Dileptons	Plasma expansion, impact parameter meter resonance melting
p_t distributions	Equation of state, Evidence of fluid flow
Strangeness and Charm	Dynamics of Expansion
Pion Correlations	Size and Lifetime of Plasma
Jets	Scattering cross section of quarks or gluons with plasma and hadronic matter

We shall discuss in detail only the electromagnetic probes and strangeness and charm in this section. The p_t distributions and hydrodynamic expansion will be discussed in later talks. The bottom line on all of these probes is that they all will involve correlations between several variables. For example, just the requirement of head-on, small impact parameter collisions requires a cut either on total multiplicity or nuclear fragmentation. Because of this often times complicated analysis of correlated variables, it is difficult to argue that any one of the probes will yield an unambiguous signal for a plasma. Nevertheless, in several cases such as photon and di-lepton probes or J/Psi production, with a little luck it may be possible to construct a convincing case that a plasma has been formed, and to measure some of its properties.

3.1 Photons and Dileptons

Quark-antiquark annihilation produces di-lepton pairs in the plasma. If we sum over all possible quark-gluon interactions in the initial and final state, then the overall rate for production of di-leptons and photons per unit time and volume is proportional to[??]

$$\frac{dN}{dtd^3xd^4q} \sim Im \int d^4x \; < J^\mu(x) J^\nu(0) > \; e^{iqx} \tag{13}$$

This assumes emission from a plasma at a fixed temperature T. The brackets <> denote a thermal expectation value. The current $J^\mu(x)$ has a real, Minkowski time argument.

There are of course a variety of non-thermal sources for di-leptons and photons. There are backgrounds for photons from π^o decays, which in the low q region obscure the signal. There may also be backgrounds for the di-leptons arising from decays of charmed particles. For large q, hard scattering processes from the initially un-thermalized beams of quarks and gluons presumably dominate. As the momentum is softened, the contributions arise from an ever more thermalized system which eventually may come from a plasma, provided backgrounds from soft hadronic decays do not become too large of a background. In this intermediate range of q, there are several thermal regions which contribute. At the higher q values, there is presumably a contribution from a quark-gluon plasma, at lower q a mixed phase of plasma and hadronic gas, and at the lowest q values larger than that for which background becomes important, there is a contribution from a hadronic gas.

To compute these distributions of photons and di-leptons, a knowledge of the space-time history of the evolution of the quark-gluon plasma isrequired.[??–??] Detailed estimates of the space-time evolution of matter produced in head-on collisions of nuclei at large A have now been carried out,[??–??] and the di-lepton distributions have been computed in detail. There has as yet been no attempt to treat non-zero impact parameter collisions. Techniques have also been developed to study the fragmentation region.[??–??] No attempt has been made to treat the pre-equilibrium region, although the cascade computation of Boal may be useful for this.[??] A treatment of the late stages in the evolution of the matter are best treated by cascade simulation of pion interactions, and again could easily be used to compute di-lepton and photon distributions.[??]

The general results of these analysis are the following:

1) For photons and di-leptons emitted from the plasma, the rapidity density of the electromagnetically produced particles is correlated with the ra-

pidity density squared of hadrons. This has been shown to be a general feature of models where the electromagnetically produced particles are produced by final state interactions of hadrons.[77] A plot of this correlation computed in a 1+1 dimensional hydrodynamic model is shown in Fig. 5.[77]

2) Pion rapidity fluctuations are correlated with fluctuations in the di-lepton and photon production rate, at the same rapidity, for thermal emission. This correlation is much different from the case for Drell-Yan pair production where there is no such correlation.

3) The rate of thermal production may be as high as 10^2 times background for not unreasonable values of the temperature. The plasma contribution is most sensitive to the values of the initial temperature when the system becomes thermalized. In Figs. 6a-6b, these thermal distributions are compared to backgrounds from Drell-Yan, and a generous estimate of backgrounds from resonances and other low p_T phenomenon. For an initial temperature of 500 Mev, the thermal signal is always 10^2 times background for masses of 2-4 Gev, as shown in Fig. 6a. For initial temperature of 240 Mev, the di-lepton spectrum is shown in Fig. 6b. Here the plasma contribution is of the same order as the Drell-Yan contribution for masses of 2-4 Gev.

4) The shape of the thermal di-lepton distribution is fairly sensitive to T_i, the largest value of the temperature for which there is a thermal distribution. The effects of a pre-equilibrium distribution of quarks and gluons has not yet been included so this conclusion is a little soft.

5) For a quark-gluon plasma at high temperature, the distribution of di-leptons is a function only of the transverse mass, $M_t = \{M^2 + p_T^2\}^{1/2}$ There should be a strong correlation between M and p_T, a correlation not present in the Drell-Yan distribution for intermediate mass pairs.

6) The distribution of di-leptons in no simple way reflects the transition temperature. This is a consequence of doing a proper 3+1 dimensional hydrodynamic computation. In 1+1 dimensional computations, the transition temperature controls the distribution in the region of $M \sim 1 - 2\ Gev$. The shape does of course weakly reflect the transition temperature, but there seems no obvious or convincing way to extract it.

7) The proposed melting of low mass resonances

such as the ρ and ω, characteristic of 1+1 dimensional hydrodynamic simulations, [??–??] is not verified in 3+1 dimensional computations. In 1+1 dimensions, the ρ and ω disappear as a resonance in the mass spectrum at large p_T since di-leptons at large p_T are emitted from a high temperature plasma. A high temperature plasma has no ρ or ω resonance. This effect disappears in the 3+1 dimensional computations because transverse expansion makes a large amount of rapidly expanding hadron gas. This transversely expanding hadron gas dominates the spectrum for masses of $M \sim 1\ Gev$ and large p_T. The melting phenomenon is presumably still effective for large mass resonances such as the J/ψ.[??]

3.2 Strange Particle Production

Strangeness has been widely suggested as a possible signal for the production of a quark-gluon plasma.[??,??] The argument for large strangeness in its most naive form follows from the observation that there are equal numbers of up, down and strange quarks in the plasma. One might naively expect that there would be roughly equal numbers of kaons and pions produced, and that the ratio of strange to non-strange baryons would be proportional to their statistical weight, $N_S/N_{NS} \sim 2/3$.

For the case of mesons, the above argument may be easily seen to be false.[??,??] In the expansion of the quark-gluon plasma, and later the hadron gas, entropy is conserved, and the pions are a result of this entropy. A better measure of the strangeness of a plasma is therefore the K/S ratio, where S is the entropy. This may be computed and shown to be smaller in a plasma than in a hadron gas for all temperatures larger than 100 Mev. The K/π ratio is therefore not a direct signal for a plasma. Further, the K/π ratio may be computed in a variety of hydrodynamic scenarios.[??–??] The result is typically $K/\pi \sim .3$. This number is a little larger than is typical of $\bar{p}p$ interactions. As has been suggested by Rafelski and Muller, perhaps only if a plasma is formed will the dynamics allow for such a large K/π ratio, and therefore is a signal of interesting dynamics, or perhaps even the production of a plasma.[??]

Strange baryons and anti-baryons may also provide a signal. Direct computations of the ratio of the ratios of strange to non-strange baryons in a plasma to that in a hadronic gas shows however that a hadronic gas is (if at all) only a little less strange than a plasma.[??,??] These estimates are done for net baryon number zero plasma, and an enhancement may exist for the plasma in the baryon number rich region. At RHIC and SPS energies, the baryon number density is effectively small at all rapidities, and this should be a good approximation. Again, although this ratio of ratios indicates a lack of a signal for equilibrium quark-gluon plasmas, the ratio of non-strange to strange baryons is large, .3-2, in either scenario for $100 Mev < T < 300 Mev$. This number is far larger than is typical of $\overline{p}p$ interactions, and again by the arguments of Rafelski and Muller, perhaps the only way to dynamically achieve this is by production of the plasma.[??] This ratio is therefore interesting for dynamical reasons.

I conclude therefore that a large strangeness signal is not a direct signal for production of a quark-gluon plasma. It is almost certainly a signal for interesting dynamics, and it may be true that the only reasonable dynamical scenarios where large strangeness may be produced involve the formation of a quark-gluon plasma.

I conclude with some very speculative remarks on the possibility of abundant soft production of charm in very high energy nuclear interactions. Estimates of the intrinsic strength of soft particle production, the string tension, suggest that this string tension might be a factor of perhaps 10-20 times the value for pp collisions, in the collisions of large nuclei such as uranium at energies in excess of 50 Gev/A center of mass energy.[??,??]. The large nuclei cause the rise in the string tension since there are a large number of overlapping strings for which the color fields add coherently. The large beam energy is required so that these multiple strings all occur at the same time, a situation which is guaranteed if the nuclei lorentz contract to a size of about .2-.3 Fermi.

If the scale for soft particle production changes by a factor of 10-20, then the intrinsic energy scale in the string tension shanges by a factor of 3-10, so which is about the ratio of charm to sraqnge quark masses. Charm in such a situation might therefore be as abundantly produced here as is the case for strangeness in pp interactions. We would have multiple charm production on an event by event basis. Such a situation would allow for a unique probe of energetic high temperature initial conditions in heavy ion collisions.

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1. Energy density scaled by T^4 as a function of T.

2. AA collision in the center of mass frame.

3. A dual parton model computation of the rapidity distribution of baryons minus anti-baryons in O-Pb collisions.

4. Flow distributions as measured by Gustafsson et. al.

5. dN/dy of hadrons scaled by $(dN/dy)^2$ of hadrons for head on AA collisions as a function of dN/dy of hadrons.

6. Di-leptons in ultra-relativistic nuclear collisions as a function of mass of di-lepton pair, (a) for an initial temperature of 500 Mev and (b) for 250 Mev.